

Exact partition function in $U(2) \times U(2)$ ABJM theory deformed by mass and Fayet-Iliopoulos terms

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Abstract

We exactly compute the partition function for $U(2)_k \times U(2)_{-k}$ ABJM theory on \mathbb{S}^3 deformed by mass m and Fayet-Iliopoulos parameter ζ . For $k = 1, 2$, the partition function has an infinite number of Lee-Yang zeros. For general k , in the decompactification limit the theory exhibits a quantum (first-order) phase transition at $m = 2\zeta$.

1 Introduction

The dynamics of two coincident M2 branes on the orbifold $\mathbb{R}^8/\mathbb{Z}_k$ is described by ABJM theory, three-dimensional $U(2)_k \times U(2)_{-k}$ supersymmetric Chern-Simons theory with bi-fundamental matter [1]. For this particular gauge group, the ABJM theory has $\mathcal{N} = 8$ superconformal symmetry and is in fact equivalent to Gustavsson-Bagger-Lambert theory [2, 3]. The partition function for the theory on \mathbb{S}^3 can be computed by supersymmetric localization [4, 5]. This theory can be deformed, preserving $\mathcal{N} = 4$ supersymmetry, by adding mass and Fayet-Iliopoulos (FI) parameters m, ζ , and the localization technique then reduces the full supersymmetric functional integral to the matrix integral [5]

$$Z = \frac{1}{4} \int \frac{d^2\mu}{(2\pi)^2} \frac{d^2\nu}{(2\pi)^2} \frac{\sinh^2 \frac{\mu_1 - \mu_2}{2} \sinh^2 \frac{\nu_1 - \nu_2}{2}}{\prod_{i,j} \cosh(\frac{\mu_i - \nu_j + m}{2}) \cosh(\frac{\mu_i - \nu_j - m}{2})} e^{\frac{ik}{4\pi} \sum_i (\mu_i^2 - \nu_i^2) - \frac{ik}{2\pi} \zeta \sum_i (\mu_i + \nu_i)} \quad (1)$$

where $i, j = 1, 2$. The parameter ζ represents a Fayet-Iliopoulos parameter for the diagonal $U(1)$ subgroup, whereas m corresponds to a mass for the chiral multiplets. The partition function should be understood as a function $Z(2\zeta, m; k)$, but for ease of presentation we will omit its arguments unless needed. For $k = 1$, the theory is mirror dual to $\mathcal{N} = 4$ supersymmetric super Yang-Mills theory with gauge group $U(2)$ coupled to a single fundamental hypermultiplet and a single adjoint hypermultiplet [5].

By shifting the integration variables, $x \equiv \mu - \zeta, y \equiv \nu + \zeta$, the partition function becomes

$$Z = \frac{1}{4} \int \frac{d^2x}{(2\pi)^2} \frac{d^2y}{(2\pi)^2} \frac{\sinh^2 \frac{x_1 - x_2}{2} \sinh^2 \frac{y_1 - y_2}{2}}{\prod_{i,j} \cosh \frac{x_i - y_j + m_1}{2} \cosh \frac{x_i - y_j - m_2}{2}} e^{\frac{ik}{4\pi} \sum_i (x_i^2 - y_i^2)}, \quad (2)$$

where m_1, m_2 are

$$m_1 = m + 2\zeta \quad \text{and} \quad m_2 = m - 2\zeta. \quad (3)$$

Note that ζ has dimension of mass. We are using units where the radius R of the three-sphere is $R = 1$.

The purpose of this note is to explicitly carry out the integration in (2). In the $m = \zeta = 0$ case, the integral was computed in [6] (a discussion of the partition function in the more general ABJ case can be found in [7]). On the other hand, the m, ζ -deformed ABJM theory was studied in [8] using the Fermi-gas formulation [9] and at large N for the $U(N)_k \times U(N)_{-k}$ gauge group in [10] (with $\zeta = 0$) and in [11] (with general $m, \zeta \neq 0$), where phase transitions in the complex parameter space generated by m_1, m_2 and $g = 2\pi i/k$ were investigated. Our explicit formula will uncover some interesting physical properties of the mass-deformed system with gauge group $U(2)_k \times U(2)_{-k}$.

The partition function (2) manifests the $m_1 \leftrightarrow m_2$ symmetry or, equivalently, $\zeta \rightarrow -\zeta$. A less obvious symmetry is $m_2 \rightarrow -m_2$, or [8, 11]

$$Z(2\zeta, m; k) = Z(m, 2\zeta; k). \quad (4)$$

For the $k = 1$ case, this symmetry already appeared in [5], where it was also explained by the fact that the corresponding brane configuration is self-mirror. The symmetry implies, in particular, that a FI-deformation ζ on the massless theory is equivalent to a mass-deformation $m = 2\zeta$ in the theory with vanishing FI-parameter. The case $m = 2\zeta$ –representing a fixed point of this symmetry– is special, as we shall shortly see. In the dual $\mathcal{N} = 4$ supersymmetric super Yang-Mills theory, $m_2 = 0$ corresponds to coupling the theory to a massless adjoint hypermultiplet.

2 Residue integration

The partition function for the m, ζ -deformed ABJM theory with $U(N)_k \times U(N)_{-k}$ gauge group can be written in the following form [5, 11]

$$Z(2\zeta, m; k) = \sum_{\rho} (-1)^{\rho} \frac{1}{N!} \int d^N \tau \frac{e^{-ikm_2 \sum_i \tau_i}}{\prod_i \cosh(k\pi\tau_i) \cosh(\pi(\tau_i - \tau_{\rho(i)}) - \frac{m_1}{2})}, \quad (5)$$

where the sum goes over permutations. The derivation uses a trigonometric identity, Fourier integrations and only holds for opposite Chern-Simons levels (see sect. 2 in [11] for details). For $N = 2$, the formula (5) then leads to the following expression

$$Z = \frac{1}{2} (Z_1 - Z_2), \quad (6)$$

with

$$Z_1 = \int d\tau_1 d\tau_2 \frac{e^{-ikm_2(\tau_1 + \tau_2)}}{\cosh(\pi k\tau_1) \cosh(\pi k\tau_2) \cosh^2\left(\frac{m_1}{2}\right)}, \quad (7)$$

and

$$Z_2 = \int d\tau_1 d\tau_2 \frac{e^{-ikm_2(\tau_1 + \tau_2)}}{\cosh(\pi k\tau_1) \cosh(\pi k\tau_2) \cosh\left(\pi(\tau_1 - \tau_2) - \frac{m_1}{2}\right) \cosh\left(\pi(\tau_1 - \tau_2) + \frac{m_1}{2}\right)}, \quad (8)$$

Using the identity

$$\frac{1}{\cosh^2 \frac{m_1}{2}} - \frac{1}{\cosh\left(\pi\tau - \frac{m_1}{2}\right) \cosh\left(\pi\tau + \frac{m_1}{2}\right)} = \frac{\operatorname{sech}^2 \frac{m_1}{2} \sinh^2 \pi\tau}{\cosh\left(\pi\tau - \frac{m_1}{2}\right) \cosh\left(\pi\tau + \frac{m_1}{2}\right)} \quad (9)$$

and the formula for the Fourier transform [11]

$$\int du \frac{e^{-ikm_2 u}}{\cosh\left(\frac{\pi k}{2}(u + v)\right) \cosh\left(\frac{\pi k}{2}(u - v)\right)} = \frac{4 \sin(km_2 v)}{k \sinh(\pi k v) \sinh m_2}, \quad (10)$$

we obtain

$$Z = \frac{1}{k^2 \sinh(m_2) \cosh^2 \frac{m_1}{2}} \int du \frac{\sin(m_2 u) \sinh^2 \frac{\pi u}{k}}{\sinh(\pi u) \cosh\left(\frac{\pi u}{k} - \frac{m_1}{2}\right) \cosh\left(\frac{\pi u}{k} + \frac{m_1}{2}\right)}. \quad (11)$$

In the limit $m_2 \rightarrow 0$, the partition function becomes

$$Z \Big|_{m_2=0} = \frac{1}{k^2 \cosh^2 \frac{m_1}{2}} \int du \frac{u \sinh^2 \frac{\pi u}{k}}{\sinh(\pi u) \cosh(\frac{\pi u}{k} - \frac{m_1}{2}) \cosh(\frac{\pi u}{k} + \frac{m_1}{2})} . \quad (12)$$

In the following, we compute the integrals (11), (12) by residue integration.

To compute (11) we follow the ideas in [6], where the partition function was computed in the case $m = \zeta = 0$.

Thus we start by writing the integrand as the product of two even functions f, g

$$Z = \frac{1}{k^2 \sinh(m_2) \cosh^2 \frac{m_1}{2}} \int du f(u) g(u) , \quad (13)$$

with

$$f(u) = \frac{\sin m_2 u}{\sinh \pi u}, \quad g(u) = \frac{\sinh^2 \frac{\pi u}{k}}{\cosh(\frac{\pi u}{k} - \frac{m_1}{2}) \cosh(\frac{\pi u}{k} + \frac{m_1}{2})} . \quad (14)$$

Under the shift $u \rightarrow u + ik$ these functions transform as

$$\begin{aligned} f(u) &\rightarrow (-)^k \cosh(m_2 k) f(u) + \text{odd function}, \\ g(u) &\rightarrow g(u) \end{aligned} \quad (15)$$

These properties imply that the integral in (13) along the curve $u = x + ik$ with $x \in \mathbb{R}$ will differ from the integration along the real axis by the factor $(-)^k \cosh(m_2 k)$. Therefore, the rectangular contour composed by the real axis, two vertical segments and the displaced real axis $u = x + ik$ becomes appropriate for residue computation in the case $m_2 \neq 0$ (see Fig.1)¹.

The residues encircled by the contour comprise the ones arising from the poles of $f(z)$ located at $z = in$ with $n = 1, \dots, k$ and those of $g(z)$ located at $z_{\pm} = \pm \frac{m_1 k}{2\pi} + i \frac{k}{2}$. The pole located at $z = ik$ does not contribute due to a double zero in the numerator of $g(z)$. Calling C the closed rectangular contour described above and $\mathcal{F}(z) = f(z)g(z)$ one finds

$$\begin{aligned} \oint_C dz \mathcal{F}(z) &= (1 - (-)^k \cosh(m_2 k)) \int du \mathcal{F}(u) \\ &= 2\pi i \left[\sum_{n=1}^{k-1} \text{Res}_{z=in} \mathcal{F}(z) + \text{Res}_{z=z_{\pm}} \mathcal{F}(z) \right] \end{aligned}$$

which gives

$$\int du \mathcal{F}(u) = \frac{2\pi i}{1 - (-)^k \cosh(m_2 k)} \left[-\frac{i}{\pi} \sum_{n=1}^{k-1} (-)^n \frac{\sin^2(\frac{n\pi}{k}) \sinh(m_2 n)}{\cosh(\frac{m_1}{2} - \frac{in\pi}{k}) \cosh(\frac{m_1}{2} + \frac{in\pi}{k})} + R_k \right] \quad (16)$$

¹It is easily seen that the vertical contours do not contribute when we push them to infinity.

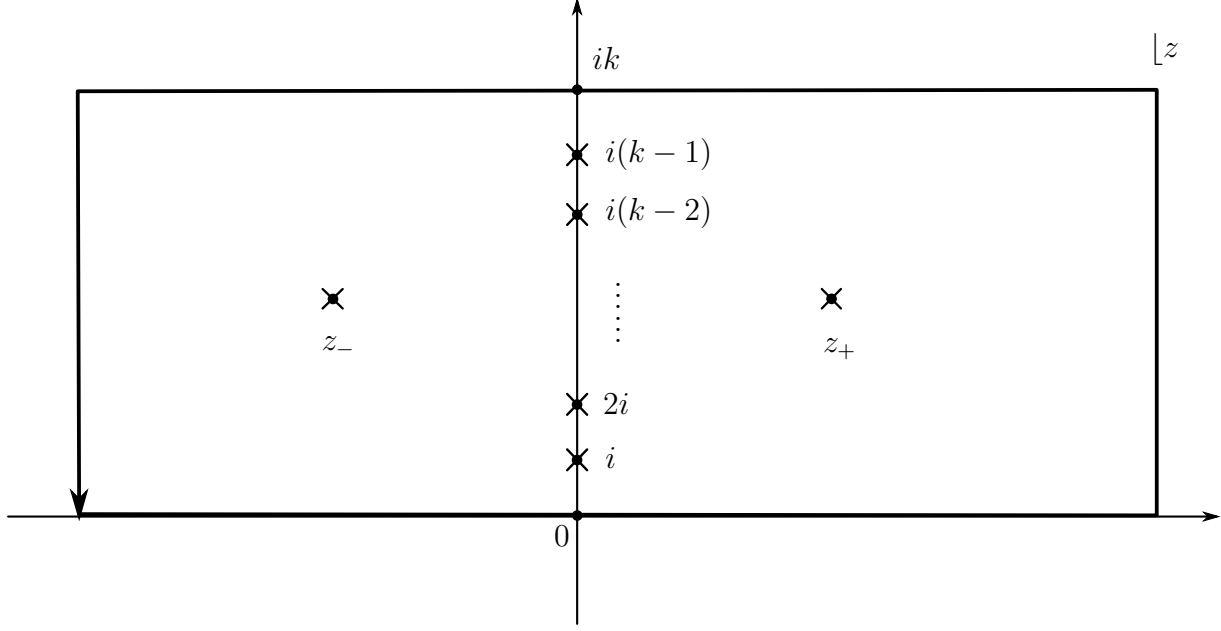


Figure 1: Rectangular contour for residue computation. The poles on the imaginary axis $z = in$ with $n = 1, \dots, k-1$ arise from the f function, while those at $z_{\pm} = \pm \frac{m_1 k}{2\pi} + i\frac{k}{2}$ follow from the g function.

where

$$R_k = \begin{cases} (-)^{\frac{k}{2}} \frac{ik}{\pi} \frac{\coth \frac{m_1}{2} \sinh \frac{km_2}{2}}{\sinh \frac{km_1}{2}} \cos \frac{km_1 m_2}{2\pi}, & k \text{ even} \\ (-)^{\frac{k+1}{2}} \frac{ik}{\pi} \frac{\coth \frac{m_1}{2} \cosh \frac{km_2}{2}}{\cosh \frac{km_1}{2}} \sin \frac{km_1 m_2}{2\pi}, & k \text{ odd} \end{cases} \quad (17)$$

Case $m_2 = 0$, k odd: it is evident from (16) that the $m_2 \rightarrow 0$ limit of (13) is smooth, the result is

$$Z \Big|_{m_2=0} = \frac{1}{k^2 \cosh^2 m} \left[\sum_{n=1}^{k-1} (-)^n \frac{n \sin^2(\frac{n\pi}{k})}{\cosh(m - \frac{in\pi}{k}) \cosh(m + \frac{in\pi}{k})} - (-)^{\frac{k+1}{2}} \frac{k^2 m \coth m}{\pi \cosh km} \right], k \text{ odd} \quad (18)$$

where we have used $m_1 = 2m$.

Case $m_2 = 0$, k even: the factor multiplying the bracket in (16) prevents taking $m_2 \rightarrow 0$ in the even k case. To compute the integral in (12) we consider

$$I = \int du \tilde{f}(u) g(u), \quad (19)$$

with $g(u)$ as in (14) and

$$\tilde{f}(u) = \frac{i}{k} \frac{(u - ik/2)^2}{\sinh \pi u}.$$

Upon integration, the odd piece in \tilde{f} vanishes against $g(u)$ and therefore the partition function (12) can be written as

$$Z \Big|_{m_2=0} = \frac{1}{k^2 \cosh^2 m} I \quad (20)$$

The shift $u \rightarrow u + ik$ in $\tilde{f}(u)$ gives

$$\tilde{f}(u) \rightarrow (-)^{k+1} \tilde{f}(-u) .$$

As discussed below (15), this property makes the rectangular contour in Fig.1 appropriate for computing I by residues.

For the residues analysis we should now consider the pole in $\tilde{f}(z)$ at the origin $z = 0$ but a zero in $g(z)$ eliminates it; along the same lines the residue from $z = ik/2$ is absent since a zero appears for \tilde{f} . Calling $\tilde{\mathcal{F}}(z) = \tilde{f}(z)g(z)$ one finds

$$\oint_C dz \tilde{\mathcal{F}}(z) = 2I,$$

on the other hand

$$\begin{aligned} \oint_C dz \tilde{\mathcal{F}}(z) &= 2\pi i \left[\sum_{n=0}^{k-1} \text{Res}_{z=in} \tilde{\mathcal{F}}(z) + \text{Res}_{z=z_{\pm}} \tilde{\mathcal{F}}(z) \right] \\ &= 2\pi i \left[\frac{i}{k\pi} \sum_{n=1}^{k-1} (-)^n \left(\frac{k}{2} - n \right)^2 \frac{\sin^2(\frac{n\pi}{k})}{\cosh(m - \frac{in\pi}{k}) \cosh(m + \frac{in\pi}{k})} + \tilde{\mathcal{R}}_k \right] \end{aligned} \quad (21)$$

where

$$\tilde{\mathcal{R}}_k = (-)^{\frac{k}{2}} \frac{2i(mk)^2 \coth(m) \sinh mk}{\pi^3 \cosh(2mk) - 1}$$

The $n = \frac{k}{2}$ term in the sum vanishes as expected. The final result is

$$\begin{aligned} Z \Big|_{m_2=0} &= -\frac{1}{k \cosh^2 m} \cdot \\ &\quad \left[\sum_{n=1}^{k-1} (-)^n \left(\frac{n}{k} - \frac{1}{2} \right)^2 \frac{\sin^2(\frac{n\pi}{k})}{\cosh(m - \frac{in\pi}{k}) \cosh(m + \frac{in\pi}{k})} + (-)^{\frac{k}{2}} \frac{2m^2 k \coth(m) \sinh mk}{\pi^2 \cosh(2mk) - 1} \right] \end{aligned} \quad (22)$$

3 Summary of results and limits

Thus we have obtained

$$Z = \frac{2}{k^2 \sinh(m_2)} \frac{1}{1 - (-1)^k \cosh(m_2 k)} (J_1 - J_2) \quad (23)$$

where

$$J_1 = \frac{1}{\cosh^2(\frac{m_1}{2})} \sum_{n=1}^{k-1} (-1)^n \frac{\sin^2(\frac{n\pi}{k}) \sinh(m_2 n)}{\cosh(\frac{m_1}{2} - \frac{in\pi}{k}) \cosh(\frac{m_1}{2} + \frac{in\pi}{k})} \quad (24)$$

and

$$J_2 = \begin{cases} (-1)^{\frac{k}{2}} \frac{2k \sinh \frac{km_2}{2}}{\sinh(m_1) \sinh \frac{km_1}{2}} \cos \frac{km_1 m_2}{2\pi}, & k \text{ even} \\ (-1)^{\frac{k+1}{2}} \frac{2k \cosh \frac{km_2}{2}}{\sinh(m_1) \cosh \frac{km_1}{2}} \sin \frac{km_1 m_2}{2\pi}, & k \text{ odd} \end{cases} \quad (25)$$

Using

$$\frac{2}{1 + \cosh \alpha} = \frac{1}{\cosh^2(\frac{\alpha}{2})}, \quad \frac{2}{1 - \cosh \alpha} = -\frac{1}{\sinh^2(\frac{\alpha}{2})}, \quad (26)$$

we can finally put the partition function in the form

$$Z \Big|_{k \text{ even}} = -\frac{1}{k^2 \sinh(m_2) \sinh^2(\frac{km_2}{2})} (J_1 - J_2) \quad (27)$$

$$Z \Big|_{k \text{ odd}} = \frac{1}{k^2 \sinh(m_2) \cosh^2(\frac{km_2}{2})} (J_1 - J_2) \quad (28)$$

In the formulas (27)-(28), the symmetry $m_1 \leftrightarrow m_2$ –which is manifest in the integral form (2)– is hidden. Interestingly, this symmetry is only recovered upon summation over n . On the other hand, the symmetry $m_2 \rightarrow -m_2$ is manifest.

Note that Z is real. While this is expected in a unitary theory, it is not generally the case in Chern-Simons theories (for a discussion, see [12]). In the present case, it is related to the fact the theory is a combination of two Chern-Simons theory with opposite levels.²

Consider, as particular examples, the important cases $k = 1, 2$. The partition functions take the form

$$Z \Big|_{k=1} = \frac{2}{\sinh(m_1) \sinh(m_2) \cosh(\frac{m_1}{2}) \cosh(\frac{m_2}{2})} \sin\left(\frac{m_1 m_2}{2\pi}\right), \quad (29)$$

$$Z \Big|_{k=2} = \frac{2}{\sinh^2(m_1) \sinh^2(m_2)} \sin^2\left(\frac{m_1 m_2}{2\pi}\right). \quad (30)$$

Now the symmetry $m_1 \leftrightarrow m_2$ has become manifest.

Note that the partition functions for $k = 1, 2$ have zeros. Restoring the R dependence, the zeros are located at

$$m_1 m_2 R^2 = 2\pi^2 n, \quad n = \pm 1, \pm 2, \dots \quad (31)$$

They represent Lee-Yang zeros (see, for example, [13]). In the infinite volume, $R \rightarrow \infty$, the zeros condense in a certain line, and a phase transition should emerge. The fact that the partition function has zeros seems to be related to the fact that the coupling,

²We thank Miguel Tierz for comments on this point.

$g = 2\pi i/k$, is imaginary for real k . Indeed, from the general expressions (24)-(25) we see that the arguments of the sine and cosine functions in (29), (30) contain a factor π/k . If the coupling g is (unphysically) continued to the real line by taking $k \rightarrow ik$, the partition function zeros would then lie on the imaginary g -axis, in accordance with the Lee-Yang theorem (see [11] for a related discussion).

For the undeformed ABJM theory, the $k = 1$ case is of special interest, since it is conjectured to describe the dynamics of two M2 branes in eleven-dimensional Minkowski spacetime. An interesting question is what is the origin of these Lee-Yang singularities in the brane realization.

The partition function $Z(2\zeta, m; k)$ does not have any zeros for $k > 2$. For higher values of k , the partition function becomes more involved, below we quote explicitly the $k = 3$ and $k = 4$ cases

$$Z \Big|_{k=3} = \frac{2}{3} \frac{2 - \sin\left(\frac{3m_1 m_2}{2\pi}\right) \operatorname{csch}\left(\frac{m_1}{2}\right) \operatorname{csch}\left(\frac{m_2}{2}\right)}{(\cosh m_1 + \cosh 2m_1)(\cosh m_2 + \cosh 2m_2)} \quad (32)$$

$$Z \Big|_{k=4} = \frac{1 - \operatorname{sech}(m_1) - \operatorname{sech}(m_2) + \cos\left(\frac{2m_1 m_2}{\pi}\right) \operatorname{sech}(m_2) \operatorname{sech}(m_1)}{8 \sinh^2 m_1 \sinh^2 m_2} \quad (33)$$

Note that the symmetry under the exchange $m_1 \leftrightarrow m_2$ is manifest.

Asymptotic formulas

Let us consider the limit of a large sphere, $mR \gg 1$, at fixed k . Assuming $m_1 > 0$, $m_2 > 0$ and restoring the R dependence, we find

$$Z \Big|_{k=1} \sim 32 e^{-\frac{3}{2}(m_1+m_2)R} \sin\left(\frac{m_1 m_2 R^2}{2\pi}\right), \quad (34)$$

$$Z \Big|_{k=2} \sim 32 e^{-2(m_1+m_2)R} \sin^2\left(\frac{m_1 m_2 R^2}{2\pi}\right), \quad (35)$$

$$Z \Big|_{k>2} \sim \frac{64}{k^2} e^{-2(m_1+m_2)R} \sin^2\left(\frac{\pi}{k}\right). \quad (36)$$

The general asymptotic formula with arbitrary sign for m_2 and $m_2 \neq 0$, is obtained by replacing m_2 by $|m_2|$.

The absolute value implies a discontinuity in the first derivative of $F = -\ln Z$. This indicates a first-order phase transition in the parameter m_2 at $m_2 = 0$, i.e., when the two mass scales $m, 2\zeta$ cross. Explicitly, at large R , we have

$$F = 2(|m_1| + |m_2|)R + O(1), \quad k > 1. \quad (37)$$

Hence

$$\frac{d\Delta F}{dm_2} \Big|_{m_2=0} = 4R, \quad \Delta F \equiv F_{m_2>0} - F_{m_2<0}. \quad (38)$$

For $k = 1$ the discontinuity in the first derivative of ΔF is equal to $3R$, as can be seen from (34).

For the general theory with gauge group $U(N)_k \times U(N)_{-k}$, large N phase transitions in the complex parameter $Ng = 2\pi iN/k$ were studied in [10, 11]. These phase transitions require taking infinite volume and, at the same time, a strong coupling limit with fixed kR – a limit that already appeared in the context of supersymmetric $U(N)$ Chern-Simons theory with massive fundamental matter in [14, 15]. It should be noted that such decompactification limit is different from the present (more physical) limit of large R at fixed k .

Another interesting aspect of (36) is that it is in a form suitable for a weak coupling expansion in powers of $1/k$:

$$Z \Big|_{k>2} \sim -\frac{32}{k^2} e^{-2(m_1+m_2)R} \sum_{n=1}^{\infty} \frac{(-1)^n}{(2n)!} \left(\frac{2\pi}{k}\right)^{2n}. \quad (39)$$

The perturbative expansion has an infinite radius of convergence. However, the original theory on the three-sphere of *finite* radius R has an asymptotic perturbative expansion, with $2n!$ asymptotic behavior for the $1/k^{2n}$ term. This can be seen by using the integral form (11) and generalizing the study of [16, 17] on the resurgence properties of the perturbation series of ABJM theory. Now, expanding the integrand in (11), one finds a series with finite radius of convergence determined by the poles of $\text{sech}(\pi u/k \pm m_1/2)$ in the complex u -plane. The integral over u then adds an extra $(2n)!$, leading to an asymptotic (but Borel summable) perturbation series.

4 The special case $m_2 = 0$

The $m_2 = 0$ case is special and must be considered separately. In particular, it represents the critical point in the phase transitions that arise in the decompactification limit. In section 2 we have obtained the following formulas:

Odd k :

$$Z \Big|_{m_2=0} = \frac{1}{k^2 \cosh^2 m} \sum_{n=1}^{k-1} (-)^n \frac{n \sin^2 \frac{\pi n}{k}}{\cosh(m + \frac{i\pi n}{k}) \cosh(m - \frac{i\pi n}{k})} + \frac{(-)^{\frac{k-1}{2}} 2m}{\pi \cosh(km) \sinh(2m)}. \quad (40)$$

Even k :

$$\begin{aligned} Z \Big|_{m_2=0} = & \frac{1}{k \cosh^2 m} \sum_{n=1}^{k-1} (-)^{n+1} \left(\frac{n}{k} - \frac{1}{2}\right)^2 \frac{\sin^2(\frac{n\pi}{k})}{\cosh(m - \frac{in\pi}{k}) \cosh(m + \frac{in\pi}{k})} \\ & + (-)^{\frac{k}{2}+1} \frac{4m^2}{\pi^2} \frac{\sinh km}{\sinh(2m)(\cosh(2mk) - 1)} \end{aligned} \quad (41)$$

In particular,

$$\begin{aligned} Z \Big|_{k=1} &= \frac{2m}{\pi \cosh(m) \sinh(2m)} , \\ Z \Big|_{k=2} &= \frac{2m^2}{\pi^2 \sinh^2(2m)} . \end{aligned} \tag{42}$$

Note that the partition function does not have zeros in this case.

Asymptotic formulas $m_2 = 0$

Consider again the limit of a large sphere, $mR \gg 1$, at fixed k , but now with $m_2 = 0$. We find

$$Z \Big|_{k=1} \sim \frac{8mR}{\pi} e^{-3mR} , \tag{43}$$

$$Z \Big|_{k=2} \sim \frac{8}{\pi^2} m^2 R^2 e^{-4mR} , \tag{44}$$

$$Z \Big|_{k>2} \sim \frac{4}{k^2} e^{-4mR} \tan^2 \frac{\pi}{k} . \tag{45}$$

Note that these formulas differ from the asymptotic formulas (34)–(36) given above for $Z(m_1, m_2)$ at $m_2 = 0$. This is expected, since the latter were obtained by assuming $|m_1 R|, |m_2 R| \rightarrow \infty$.

Unlike the $m_2 \neq 0$ case, the perturbation series for this flat-theory limit has now finite radius of convergence $|\pi/k| < \pi/2$, therefore perturbation series is convergent for all $k > 2$, where the formula applies. On the other hand, just like the general $m_2 \neq 0$ case, the theory on a finite-radius \mathbb{S}^3 has an asymptotic perturbation series with $2n!$ asymptotic behavior.

Finally, it would be interesting to study supersymmetric Wilson loops in the present mass/FI deformed theory, along the lines of [18].

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References

- [1] O. Aharony, O. Bergman, D. L. Jafferis and J. Maldacena, “N=6 superconformal Chern-Simons-matter theories, M2-branes and their gravity duals,” JHEP **0810**, 091 (2008) [arXiv:0806.1218 [hep-th]].

- [2] A. Gustavsson, “Algebraic structures on parallel M2-branes,” Nucl. Phys. B **811**, 66 (2009) [arXiv:0709.1260 [hep-th]].
- [3] J. Bagger and N. Lambert, “Gauge symmetry and supersymmetry of multiple M2-branes,” Phys. Rev. D **77**, 065008 (2008) [arXiv:0711.0955 [hep-th]].
- [4] A. Kapustin, B. Willett and I. Yaakov, “Exact Results for Wilson Loops in Superconformal Chern-Simons Theories with Matter,” JHEP **1003**, 089 (2010) [arXiv:0909.4559 [hep-th]].
- [5] A. Kapustin, B. Willett and I. Yaakov, “Nonperturbative Tests of Three-Dimensional Dualities,” JHEP **1010**, 013 (2010) [arXiv:1003.5694 [hep-th]].
- [6] K. Okuyama, “A Note on the Partition Function of ABJM theory on S^3 ,” Prog. Theor. Phys. **127**, 229 (2012) [arXiv:1110.3555 [hep-th]].
- [7] H. Awata, S. Hirano and M. Shigemori, “The Partition Function of ABJ Theory,” PTEP **2013**, 053B04 (2013) [arXiv:1212.2966].
- [8] N. Drukker and J. Felix, “3d mirror symmetry as a canonical transformation,” JHEP **1505**, 004 (2015) [arXiv:1501.02268 [hep-th]].
- [9] M. Marino and P. Putrov, “ABJM theory as a Fermi gas,” J. Stat. Mech. **1203**, P03001 (2012) [arXiv:1110.4066 [hep-th]].
- [10] L. Anderson and K. Zarembo, “Quantum Phase Transitions in Mass-Deformed ABJM Matrix Model,” JHEP **1409**, 021 (2014) [arXiv:1406.3366 [hep-th]].
- [11] L. Anderson and J. G. Russo, “ABJM Theory with mass and FI deformations and Quantum Phase Transitions,” JHEP **1505**, 064 (2015) [arXiv:1502.06828 [hep-th]].
- [12] C. Closset, T. T. Dumitrescu, G. Festuccia, Z. Komargodski and N. Seiberg, “Contact Terms, Unitarity, and F-Maximization in Three-Dimensional Superconformal Theories,” JHEP **1210**, 053 (2012) [arXiv:1205.4142 [hep-th]].
- [13] C. Itzykson and J. M. Drouffe, “Statistical Field Theory. Vol. 1: From Brownian Motion To Renormalization And Lattice Gauge Theory,” Cambridge, UK: Univ. Pr. (1989) 1-403
- [14] A. Barranco and J. G. Russo, “Large N phase transitions in supersymmetric Chern-Simons theory with massive matter,” JHEP **1403**, 012 (2014) [arXiv:1401.3672 [hep-th]].
- [15] J. G. Russo, G. A. Silva and M. Tierz, “Supersymmetric U(N) Chern-Simons Matter Theory and Phase Transitions,” Commun. Math. Phys. **338**, no. 3, 1411 (2015) [arXiv:1407.4794 [hep-th]].

- [16] J. G. Russo, “A Note on perturbation series in supersymmetric gauge theories,” JHEP **1206**, 038 (2012) [arXiv:1203.5061 [hep-th]].
- [17] I. Aniceto, J. G. Russo and R. Schiappa, “Resurgent Analysis of Localizable Observables in Supersymmetric Gauge Theories,” JHEP **1503**, 172 (2015) [arXiv:1410.5834 [hep-th]].
- [18] S. Hirano, K. Nii and M. Shigemori, “ABJ Wilson loops and Seiberg duality,” PTEP **2014**, no. 11, 113B04 (2014) [arXiv:1406.4141 [hep-th]].